

## PLANET FORMATION IN SMALL SEPARATION BINARIES: NOT SO EXCITED AFTER ALL

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## ABSTRACT

Existence of planets in binaries with relatively small separations (around 20 AU), such as  $\alpha$  Centauri or  $\gamma$  Cephei poses severe challenges to standard planet formation theories. The problem lies in the vigorous secular excitation of planetesimal eccentricities at separations of several AU, where some of the planets are found, by the massive, eccentric stellar companions. High relative velocities of planetesimals preclude their growth in mutual collisions for a wide range of sizes, from below 1 km up to several hundred km, resulting in fragmentation barrier to planet formation. Here we show that rapid apsidal precession of planetesimal orbits, caused by the gravity of the circumstellar protoplanetary disk, acts to strongly reduce eccentricity excitation, lowering planetesimal velocities by an order of magnitude or even more at 1 AU. By examining the details of planetesimal dynamics we demonstrate that this effect eliminates fragmentation barrier for in-situ growth of planetesimals as small as  $\lesssim 10$  km even at separations as wide as 2.6 AU (semi-major axis of the giant planet in HD 196885), provided that the circumstellar protoplanetary disk is relatively massive,  $\sim 0.1M_{\odot}$ .

*Subject headings:* planets and satellites: formation — protoplanetary disks — planetary systems — binaries: close

## 1. INTRODUCTION.

About 20% of planets detected via stellar radial velocity variations reside in binaries (Desidera & Barbieri 2007). The majority of these systems are wide separation binaries, with semi-major axis  $a_b \gtrsim 30$  AU. At the same time, four relatively small separation binaries with  $a_b \approx 20$  AU (HD 196885,  $\gamma$  Cephei, Gl 86 and HD 41004; Chauvin et al. 2011) are also known to harbor giant planets with projected masses  $M_{pl} \sin i \approx (1.6 - 4)M_J$ . In these systems the mass of the secondary star (we call “secondary” the binary component other than the star orbited by the planet, which we denote as “primary”)  $M_s$  is found to be close to  $0.4M_{\odot}$  and binary eccentricity  $e_b$  is close to 0.4. Also, Dumusque et al. (2012) have recently announced an Earth-mass companion to  $\alpha$  Centauri B, a member of the binary (or, possibly, a triple) with  $a_b = 17.6$  AU,  $e_b = 0.52$ , and  $M_s = 1.1M_{\odot}$ . In this system planet orbits the star at  $\approx 0.04$  AU separation.

The uniqueness of these systems lies in the fact that forming planets in them is known to provide extreme challenge to planet formation theories (Zhou et al. 2012). With the exception of  $\alpha$  Cen and Gl 86, planets in these binaries reside in rather wide orbits, with planetary semi-major axes  $a_{pl} \approx 1.6 - 2.6$  AU. In-situ formation of these gas giants is expected to proceed through continuous agglomeration of planetesimals at these locations, starting from very small objects (easily  $\lesssim 1$  km). However, gravitational perturbations from the eccentric stellar companion inevitably result in rapid secular evolution (Heppenheimer 1978), driving planetesimal eccentricities far above the level at which bodies can avoid destruction in mutual collisions (Th  bault et al. 2008). This problem, which is often called *collisional* or *fragmentation barrier*, is especially severe for small planetesimals,  $1 - 10^2$  km in size, for which the ratio of binding to kinetic energy is small. It is also more pronounced far from the primary,

where the secular forcing by the companion is strongest and planetesimal eccentricities are high.

Marzari & Scholl (2000) suggested that a combination of secular forcing by the companion and gas drag acting on small ( $1 - 10$  km) planetesimals leads to apsidal alignment of their orbits, resulting in smaller relative velocities, and allowing colliding objects to grow. However, Th  bault et al. (2006, 2008) demonstrated that the planetesimal size-dependence of apsidal alignment acts to break orbital phasing between objects of different sizes, resulting in high velocity collisions between them and reinforcing collisional barrier.

Interestingly, most studies of planetesimal growth in small separation binaries have included the effect of the protoplanetary disk on planetesimal dynamics only via associated gas drag (Th  bault et al. 2004, 2006, 2008, 2009; Paardekooper et al. 2008; Paardekooper & Leinhardt 2010), without accounting for the *gravitational* effect of the disk. Batygin et al. (2011) have considered disk gravity in the context of planet formation and evolution in systems with highly misaligned, distant ( $10^2 - 10^3$  AU) stellar companions, affected by the Lidov-Kozai effect (Lidov 1961; Kozai 1962). However, this effect is probably irrelevant for planetesimal dynamics in small separation (tens of AU) binaries, which are likely coplanar with circumstellar disks.

In this Letter we show that apsidal precession of planetesimal orbits induced by disk gravity dominates secular evolution of planetesimals at separations of several AU. As a result, relative velocities at which bodies collide are reduced, sometimes by more than an order of magnitude. In massive disks this effect presents a natural solution of the fragmentation barrier issue for the in-situ formation of the giant planets in small separation binaries, such as  $\gamma$  Cephei.

## 2. SECULAR EVOLUTION.

We consider planetesimal motion as Keplerian motion around the primary perturbed by the gravity of the

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companion, that moves on larger, eccentric orbit, and the disk. Mass of the primary is  $M_p$ , and we define  $\mu \equiv M_s/(M_p + M_s)$ . We assume eccentricity of the stellar binary  $e_b$  to be small and planetesimal orbit to be *coplanar* with the binary. Planetesimals are immersed in a massive, axisymmetric gaseous disk, characterized by surface density  $\Sigma(r)$  specified in §2.2.1.

Assuming  $e \ll 1$  the secular (averaged over the planetesimal and binary orbital motion) disturbing function for a planetesimal with semimajor axis  $a$  and eccentricity vector  $\mathbf{e} = (k, h) = (e \cos \varpi, e \sin \varpi)$  (with apsidal angle  $\varpi$  counted from the binary apsidal line, which is assumed fixed<sup>2</sup>) is (Murray & Dermott 1999)

$$R = na^2 \times \left[ \frac{1}{2} (A + \dot{\varpi}_d) (h^2 + k^2) - Bk \right], \quad (1)$$

where

$$A = \frac{3}{4} \mu \frac{n_b^2}{n} \left( 1 + \frac{3}{2} e_b^2 \right), \quad (2)$$

$$B = \frac{15}{16} e_b \mu \frac{n_b^2}{n} \frac{a}{a_b} \left( 1 + \frac{5}{2} e_b^2 \right), \quad (3)$$

and

$$\dot{\varpi}_d = -\frac{1}{2n} \left( \frac{2}{r} \frac{\partial U_d}{\partial r} + \frac{\partial^2 U_d}{\partial r^2} \right) \Big|_{r=a} \quad (4)$$

is the precession frequency of planetesimal orbit due to the disk potential  $U_d$ . Here  $n_b = [G(M_p + M_s)/a_b^3]^{1/2}$  and  $n = (GM_p/a^3)^{1/2}$  are the mean motions of the binary and planetesimal, respectively. The contribution to  $R$  proportional to  $\dot{\varpi}_d$  arises from expansion of the disk potential along the eccentric planetesimal orbit and averaging over its mean longitude.

Evolution equations for  $h$  and  $k$  are written using  $dh/dt = (na^2)^{-1} \partial R / \partial k$ ,  $dk/dt = -(na^2)^{-1} \partial R / \partial h$  as

$$\frac{dh}{dt} = (A + \dot{\varpi}_d)k - B, \quad \frac{dk}{dt} = -(A + \dot{\varpi}_d)h. \quad (5)$$

These equations agree with the work of Marzari & Scholl (2000) as long as  $\dot{\varpi}_d = 0$ .

We write down the solution for  $\mathbf{e}(t) = \mathbf{e}_{\text{free}}(t) + \mathbf{e}_{\text{forced}}(t)$ , where

$$\begin{Bmatrix} k_{\text{free}}(t) \\ h_{\text{free}}(t) \end{Bmatrix} = e_{\text{free}} \begin{Bmatrix} \cos[(A + \dot{\varpi}_d)t + \varpi_0] \\ \sin[(A + \dot{\varpi}_d)t + \varpi_0] \end{Bmatrix}, \quad (6)$$

$\varpi_0$  is a constant, and

$$\begin{Bmatrix} k_{\text{forced}}(t) \\ h_{\text{forced}}(t) \end{Bmatrix} = e_{\text{forced}} \begin{Bmatrix} 1 \\ 0 \end{Bmatrix}, \quad e_{\text{forced}} = \frac{B}{A + \dot{\varpi}_d}. \quad (7)$$

Thus, free eccentricity vector  $\mathbf{e}_{\text{free}}$  rotates at a rate  $A + \dot{\varpi}_d$  around the endpoint of the *fixed* forced eccentricity vector  $\mathbf{e}_{\text{forced}}$ . Note that setting  $\dot{\varpi}_d = 0$  we reproduce the solution obtained by Heppenheimer (1978).

Planetesimals starting on circular orbits have  $e_{\text{free}} = e_{\text{forced}}$  so that their eccentricity oscillates with amplitude  $2e_{\text{forced}}$  and period  $2\pi/(A + \dot{\varpi}_d)$ .

<sup>2</sup> Precession period of the secondary orbit exceeds important timescales of the problem.

## 2.1. Disk model.

We model the disk as a constant  $\dot{M}$  disk extending out to the outer truncation radius  $r_t$ . Numerical simulations of accretion disks in binaries suggest that  $r_t \sim (0.2 - 0.4)a_b$  (Zhou et al. 2012), depending on  $e_b$  and  $\mu$ . In our study we will commonly take  $r_t = 0.25a_b$ .

Constant  $\dot{M}$  assumption is a necessary simplification, which ignores the details of the disk structure at  $r \sim r_t$ . Assuming viscosity  $\nu$  in the disk to be well described by the radius-independent effective  $\alpha$ -parameter (Shakura & Sunyaev 1973), one can write

$$\Sigma(r) = \frac{\dot{M}}{3\pi\nu} = \frac{\Omega\dot{M}}{3\pi\alpha c_s^2}. \quad (8)$$

From this it is obvious that if the disk temperature scales as  $T \propto r^{-q}$  then the surface density behaves as

$$\Sigma(r) = \Sigma_0 \left( \frac{r_0}{r} \right)^p, \quad p = \frac{3}{2} - q, \quad (9)$$

where  $r_0$  is some fiducial radius and  $\Sigma_0 = \Sigma(r_0)$ .

Models of protoplanetary disks typically find  $q$  to be close to  $1/2$ , so that  $p \approx 1$ . In particular, the passive disk model of Chiang & Goldreich (1997) has  $q = 3/7$ , so that  $p - 1 = 1/14$ . Outer parts of a disk in a binary are additionally heated by the radiation of the companion and tidal dissipation, so that  $q$  may be lower than  $3/7$  even for a passive disk. For simplicity, in our calculations we will assume  $p = 1$ , which corresponds to a classical Mestel disk (Mestel 1963) if  $r_t \rightarrow \infty$ .

Assuming that the disk has power law profile (9) with  $p = 1$  all the way to  $r_t$  we can express its surface density via the total disk mass  $M_d$  enclosed within  $r_t$  as

$$\Sigma(r) \approx \frac{M_d}{2\pi r_t} r^{-1} \approx 2800 \text{ g cm}^{-2} \frac{M_d}{10^{-2} M_\odot} r_{t,5}^{-1} r_1^{-1}, \quad (10)$$

where  $r_{t,5} \equiv r_t/(5\text{AU})$ , and  $r_1 \equiv r/(1\text{AU})$ . Interestingly, gas surface density at 1 AU in such a disk with  $M_d = 0.01M_\odot$  is not very different from that in a Minimum Mass Solar Nebula (MMSN; Hayashi 1981).

## 2.2. Precession due to the disk.

A disk with the density profile (9) with  $p = 1$  extending to infinity is known to have constant circular velocity (Mestel 1963)

$$v_c = \left( r \frac{\partial U_d}{\partial r} \right)^{1/2} = (2\pi G \Sigma_0 r_0)^{1/2}. \quad (11)$$

Expressing  $\partial U_d / \partial r$  from this relation and plugging it in equation (4) we find  $\dot{\varpi}_d = -\pi G \Sigma(r)/(nr)$ , where from now on we use  $r$  instead of  $a$ . Even though the circumprimary disk in our problem is truncated at  $r_t$  this expression should still be able to give us a reasonable estimate of the precession rate  $\dot{\varpi}_d$  due to the disk potential for  $r \lesssim r_t$ . Using equation (10) we find

$$\dot{\varpi}_d \approx -\frac{GM_d}{nr_t} r^{-2} = -n \frac{M_d}{M_p} \frac{r}{r_t}. \quad (12)$$

Note that  $\dot{\varpi}_d$  varies rather weakly with  $r$ , as  $r^{-1/2}$ , which is consistent with Batygin et al. (2011).

### 2.3. Planetesimal eccentricities.

To assess the role of disk-driven precession on secular evolution of planetesimals we evaluate

$$\frac{|\dot{\varpi}_d|}{A} \approx \frac{4}{3} \frac{M_d}{M_s} \frac{a_b^3}{r_t r^2} \approx 20 \frac{M_d/M_s}{10^{-2}} \frac{a_{b,20}^3}{r_{t,5}} r_1^{-2}, \quad (13)$$

where  $a_{b,20} \equiv a_b/(20\text{AU})$ . In making this estimate we neglected the term quadratic in  $e_b$  in equation (2). It is obvious that dynamics of planetesimals at several AU is strongly affected by the disk-driven precession. Indeed,  $|\dot{\varpi}_d|$  exceeds  $A$  for

$$r \lesssim r_{cr} \approx 4.6\text{AU} \left( \frac{M_d/M_s}{10^{-2}} \frac{a_{b,20}^3}{r_{t,5}} \right)^{1/2}, \quad (14)$$

i.e. over essentially the whole assumed extent of the disk even for the disk mass as small as  $\sim 10^{-2} M_\odot$ . Thus, if we are interested in planet formation at 2–3 AU we can neglect planetesimal precession due to the secondary compared to the disk-driven precession, i.e. neglect  $A$  compared to  $\dot{\varpi}_d$  in equation (7) and other formulae.

Equation (7) then predicts that the amplitude of eccentricity oscillations is

$$e^{\text{disk}}(r) = \frac{2B}{|\dot{\varpi}_d|} \approx \frac{15}{8} e_b \frac{M_s}{M_d} \frac{r^3 r_t}{a_b^4} \quad (15)$$

$$\approx 3 \times 10^{-3} \frac{e_b}{0.5} \frac{0.01}{M_d/M_s} \frac{r_{t,5}}{a_{b,20}^4} r_1^3, \quad (16)$$

where we again neglected  $e_b^2$  term in equation (3). This is to be compared with

$$e^{\text{n/disk}}(r) = \frac{5}{2} \frac{r}{a_b} e_b \approx 6.3 \times 10^{-2} \frac{e_b}{0.5} \frac{r_1}{a_{b,20}}, \quad (17)$$

which one finds neglecting disk-driven precession, i.e. dropping  $\dot{\varpi}_d$  in equation (7). It is obvious that neglecting disk-driven precession leads to an *overestimate* of the planetesimal eccentricity at  $\sim \text{AU}$  separations by more than an order of magnitude. This has important consequences for planetesimal growth as we discuss further.

### 2.4. Gas drag.

Equation (15) accounts for the presence of the disk only through the precession caused by its gravity. However, for small planetesimals the effect of gas drag is also important. Assuming quadratic drag force in the form  $\mathbf{F} \approx -v \rho_g / (\rho d)$  (here  $d$  and  $\rho$  are the object's radius and bulk density,  $\rho_g \approx \Sigma/h$  is the gas density and  $h$  is the disk scale height) we account for its effect on planetesimal dynamics by adding terms  $-D\{h, k\} (h^2 + k^2)^{1/2}$  to the first and second equations (5), respectively (Marzari & Scholl 2000). Here

$$D = n \frac{\rho_g r}{\rho d} = n \frac{\Sigma}{\rho d} \frac{r}{h}. \quad (18)$$

For small planetesimal sizes (to be specified later by equation (20)), in the gas drag-dominated regime, the drag force balances eccentricity excitation due to the secondary, i.e. the  $B$  term in the first equation (5). This results in the following estimate for the gas drag-mediated planetesimal eccentricity:

$$e_{gas} \approx \left( \frac{B}{D} \right)^{1/2} = \left( e_b \frac{M_s}{M_p} \frac{h}{r} \frac{\rho d}{\Sigma} \right)^{1/2} \left( \frac{r}{a_b} \right)^2. \quad (19)$$

This expression agrees with Paardekooper et al. (2008) and predicts that  $e_{gas} \propto d^{1/2}$ . As a result, for small bodies one finds  $e_{gas} < e^{\text{disk}}$ .

The transition between the drag-dominated behavior (19) and the drag-free eccentricity scaling (15) occurs at the planetesimal size  $d_{gas}$  where these two equations yield the same eccentricity:

$$d_{gas} \approx \frac{4n}{B} \frac{r}{h} \frac{\Sigma}{\rho} e_{\text{forced}}^2 = \frac{15}{8\pi} e_b \frac{r}{h} \frac{M_s}{M_d} \frac{M_p r_t r}{\rho a_b^4} \quad (20)$$

$$\approx 1 \text{ km} \frac{e_b}{0.5} \frac{r/h}{30} \frac{0.01}{M_d/M_s} \frac{M_{p,1} r_{t,5} r_1}{\rho_3 a_{b,20}^4}, \quad (21)$$

where  $\rho_3 \equiv \rho/(3 \text{ g cm}^{-3})$ ,  $M_{p,1} \equiv M_p/M_\odot$  and we have used equation (10).

Planetesimal eccentricity behaves according to formula (19) for  $d \lesssim d_{gas}$ , and switches to drag-free regime (15) for  $d \gtrsim d_{gas}$ , see Figure 1. The dependence of  $d_{gas}$  on gas disk density and mass —  $d_{gas} \propto M_d^{-1}$  — is somewhat counter-intuitive, since higher gas density results in stronger drag, making it more important for *larger* bodies. However, the gas-free planetesimal eccentricity (15) is itself a function of  $M_d$  and decreases *faster* with increasing  $M_d$  than does  $e_{gas}$ , explaining the nontrivial  $d_{gas}(M_d)$  dependence.

### 3. IMPLICATIONS FOR PLANET FORMATION.

Planetesimals grow in mutual collisions as long as their encounter velocity  $v_{coll}$  (measured at infinity) is such that collisions do not result in the net loss of mass. The conditions for this depend, in particular, on planetesimal size and on whether planetesimals are strength- or gravity-dominated. Using results of Leinhardt & Stewart (2012) for collisions of equal mass (the most disruptive) strengthless bodies we roughly estimate the condition for planetesimal growth to be

$$v_{coll} \lesssim 2v_{esc}, \quad (22)$$

where the escape speed from the surface of an object of radius  $d$  and bulk density  $\rho$  is

$$v_{esc} = \left( \frac{8\pi}{3} G \rho \right)^{1/2} d \approx 1.3 \text{ m s}^{-1} \rho_3^{1/2} d_1 \quad (23)$$

(here  $d_1 \equiv d/(1\text{km})$ ). It becomes harder to break planetesimals when they are small enough for their internal strength to dominate over the gravitational energy, which is expected to happen for  $d \lesssim d_s \sim 10 \text{ km}$  (Holsapple 1994).

Planetesimal collisions occur at velocity of order  $v_{coll}(r) \approx e^{\text{disk}} v_K$ , where  $v_K = nr$  is the Keplerian speed. Using expression (15) we find

$$v_{coll}(r) \approx 90 \text{ m s}^{-1} \frac{e_b}{0.5} \frac{0.01}{M_d/M_s} \frac{M_{p,1}^{1/2} r_{t,5}}{a_{b,20}^4} r_1^{5/2}. \quad (24)$$

Plugging equations (23) and (24) into the condition (22) we find that erosion in equal-mass planetesimal collisions is avoided for bodies with  $d \gtrsim d_{coll}$ , where

$$d_{coll} \approx 35 \text{ km} \frac{e_b}{0.5} \frac{0.01}{M_d/M_s} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} \frac{r_{t,5}}{a_{b,20}^4} r_1^{5/2}. \quad (25)$$

Thus, for the fiducial binary parameters adopted here and for  $M_d \sim 10^{-2} M_\odot$  only planetesimals larger than

$\approx 35$  km would be able to grow at 1 AU. At 2 AU — the semi-major axis of  $\gamma$  Cephei Ab — only bodies larger than 200 km in radius would be able to survive in equal-mass collisions.

However, in the absence of a disk the problem is much worse: evaluating collisional velocity as  $v_{coll} = e^{n/disk} v_K$  using equation (17) and applying condition (22) one finds that in the absence of disk-induced precession only planetesimals larger than

$$d_{coll}^{n/disk} \approx 700 \text{ km} \frac{e_b}{0.5} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} a_{b,20}^{-1} r_1^{1/2} \quad (26)$$

are able to survive in equal-mass collisions. Clearly, collisional barrier appears *far* more severe if one disregards the effects of the disk on the secular evolution of planetesimals. We compare the behavior of  $d_{gas}$ ,  $d_{coll}$ , and  $d_{coll}^{n/disk}$  as a function of  $r$  in Figure 1.

We also point out that  $d_{coll}$  is very sensitive to the binary semi-major axis  $a_b$ , unlike  $d_{coll}^{n/disk}$ , see equations (25) and (26): increasing  $a_b$  from 20 AU to 30 AU reduces  $d_{coll}$  by a factor of 5. The size of the region where disk-driven precession dominates secular evolution also expands rapidly with increasing  $a_b$ , see equation (14). To summarize, properly accounting for the disk gravity considerably alleviates the collisional barrier in binaries, certainly for  $r \lesssim 1$  AU.

#### 4. PLANETESIMAL ACCRETION IS POSSIBLE IN MASSIVE DISKS

We now propose a solution to the problem of planetesimal accumulation in binaries, raised in §1. We argue that if

- disk is *massive*,  $M_d \sim 0.1 M_\odot$ ,
- planetesimals are *strength-dominated* below  $\sim 10$  km,

then the fragmentation barrier can be overcome even at separations of  $\approx 2$  AU, where planets in several binaries are found.

Equation (25) shows that higher  $M_d$  results in smaller planetesimal size  $d_{coll}$ , below which strengthless objects are destroyed or eroded in equal-mass collisions. Smaller planetesimals are (1) more resistive to collisional erosion because of their internal strength and (2) stronger affected by gas drag. When the latter dominates, planetesimal velocities are reduced and collisions are less destructive. However, increasing  $M_d$  reduces not only  $d_{coll}$  but also  $d_{gas}$  in such a way that  $d_{coll}/d_{gas}$  stays constant. At 1 AU this ratio is about 30 so that independent of  $M_d$  there is still a significant “danger zone” between the planetesimal size  $d_{gas}$  below which gas drag lowers  $v_{coll}$  helping accretion and the radius  $d_{coll}$  above which colliding strengthless bodies can grow, see Figure 1.

On the other hand, equation (26) predicts that  $d_{coll} \lesssim d_s$  at  $r = 2$  AU for  $M_d/M_s = 0.2$ , if internal strength dominates over the gravitational binding energy of the body with  $d_s = 10$  km making it harder to erode or destroy. This is a resolution of the collisional barrier problem in binaries that we favor in this work.

To avoid fragmentation barrier we thus require that  $d_{coll} < d_s$ . Using equation (25) we can rephrase this condition in the form of a lower limit on the disk mass

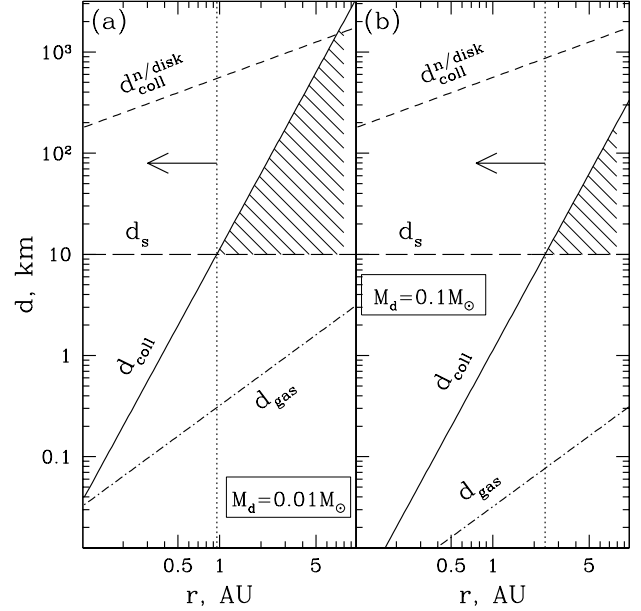


FIG. 1.— Characteristic planetesimal sizes vs. radius for two disk masses  $M_d$ : (a)  $0.01 M_\odot$  and (b)  $0.1 M_\odot$ . We display  $d_{coll}$  (solid, eq. [25]),  $d_{gas}$  (dot-dashed, eq. [20]),  $d_{coll}^{n/disk}$  (short-dashed, eq. [26]), and planetesimal radius  $d_s = 10$  km below which we consider objects as strength-dominated (long-dashed). The two latter sizes do not depend on  $M_d$ . Calculations are done for  $e_b = 0.4$ ,  $M_s = 0.4 M_\odot$ ,  $a_b = 20$  AU,  $M_p = M_\odot$  (typical for small separation binaries, Chauvin et al. 2011),  $r_t = 5$  AU, and  $r/h = 30$ . Planetesimals in the shaded region (“danger zone”) get destroyed in equal-mass collisions according to criterion (22) precluding planetary growth at corresponding separations. Accretion-friendly zone is to the left of the vertical dotted line in each plot; it is wider for higher  $M_d$  and extends to  $\approx 2.5$  AU for  $M_d = 0.1 M_\odot$ .

at a given separation  $a_{pl}$  from the primary:

$$\frac{M_d}{M_s} \gtrsim 0.035 \frac{e_b}{0.5} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} \frac{r_t/a_b}{0.25} a_{pl}^{5/2} a_{b,20}^{-3}. \quad (27)$$

In Figure 2 we illustrate this constraint as a function of the binary semi-major axis, for different values of  $a_{pl}$ . It is clear that in *very small* separation binaries with  $a_b = 10$  AU growing planets even at 1 AU requires a massive disk,  $M_p \approx 0.2 M_s$ .

External companions in giant planet-hosting binaries typically have mass  $M_s \approx 0.4 M_\odot$  (Chauvin et al. 2011), meaning that our scenario of planet formation at 2 AU needs  $M_d \approx 0.1 M_\odot$ . Such disk mass may seem high but it is also the case that planet-hosting binary systems with  $a_b \approx 20$  AU that we consider here contain more mass in total than the descendants of the typical T Tauri stars.

One might worry that such massive disks would be prone to gravitational instability (GI). With the density profile (10) we estimate the Toomre  $Q \equiv nc_s/(\pi G \Sigma)$  ( $c_s$  is the sound speed) as

$$Q \approx 2 \frac{M_p}{M_d} \frac{h}{r} \frac{r_t}{r} \approx 3 \left( \frac{0.1}{M_d/M_p} \right) \left( \frac{30}{r/h} \right) \frac{r_{t,5}}{r_1}. \quad (28)$$

Thus, even for  $M_d = 0.1 M_p \approx 0.1 M_\odot$  the disk is at most marginally unstable to GI at 2 AU. However, even if it were unstable, the surface density and optical depth at this distance would be so high that the cooling time

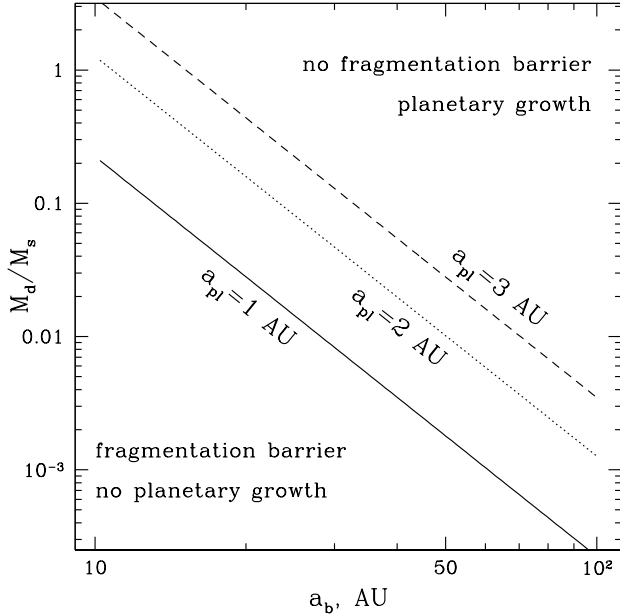


FIG. 2.— Plot of  $a_b - M_d/M_s$  phase space illustrating conditions under which planets can form in binaries, assuming  $d_s = 10$  km. Different curves show the relation (27) for different values of the planetary semi-major axis  $a_{pl}$ , interior to which planet formation is possible:  $a_{pl} = 1$  AU (solid),  $a_{pl} = 2$  AU (dotted),  $a_{pl} = 3$  AU (dashed). Calculation assumes  $e_b = 0.4$ , fixed  $r_t/a_b = 0.25$ ,  $a_b = 20$  AU, and  $M_p = M_\odot$ . At a given  $a_{pl}$  fragmentation barrier is avoided and planet formation proceeds smoothly through the planetesimal stage to the right (and above) of the corresponding line.

would far exceed the local dynamical time, making planet formation by direct disk fragmentation impossible (Gammie 2001; Rafikov 2005). Instead, the disk would slowly evolve under the action of gravitoturbulence (Rafikov 2009).

On the other hand, high  $M_d$  simplifies planet formation in other ways. In particular, planets in these systems are quite massive (Chauvin et al. 2011) and larger  $M_d$  provides mass reservoir for their assembly. Higher surface density of the protoplanetary disk also means larger isolation mass (Lissauer 1993) possibly making it high enough at 2 AU to trigger core accretion without the need to go through the long-lasting stage of giant impacts (Chambers 2004). Higher  $\Sigma$  likely implies larger dead zone (Gammie 1996) in the disk, providing quiet conditions for planetesimal formation and growth, and resulting in smaller viscosity, which, possibly, means longer

disk lifetime. The timescale on which planets form also goes down as  $M_d$  increases.

Another potential solution to the collisional barrier problem in binaries is the direct formation of large planetesimals by e.g. streaming and/or gravitational instabilities (Johansen et al. 2007; Thébault 2011). Large  $M_d$  and small  $d_{coll}$  are helpful for this mechanism as well since to overcome fragmentation barrier in a massive disk such instabilities would only need to produce bodies with sizes of tens of km, rather than  $\sim 10^3$  km dwarf planets.

## 5. DISCUSSION.

We now mention several additional factors that may strengthen or weaken our conclusions. First, our collisional growth condition (22) may be too stringent. Previously, using a more refined fragmentation criterion Thébault (2011) found that in HD 196885 ( $M_p = 1.3M_\odot$ ,  $a_b = 21$  AU,  $e_b = 0.42$ ) planetesimal growth is possible even in the absence of disk-driven precession at 2.6 AU as long as the planetesimal size exceeds 250 km. However, according to our formula (26) with the same assumptions (same system parameters and  $\dot{\omega}_d = 0$ ), growth is possible only for  $d \gtrsim 10^3$  km. Thus, our fragmentation criterion likely *overestimates* planetesimal size above which objects grow efficiently, and, in fact, it might be *easier* to overcome the fragmentation barrier with the more realistic growth condition than our simple criterion (22). At a given distance this would *lower* the value of  $M_d$  needed to overcome fragmentation barrier.

Second, even if planetesimals are collisionally weak their growth may still proceed mainly via unequal-mass collisions (which more frequently result in mergers) if the number of relatively massive objects is small (Thébault 2011). Outward migration of planets by scattering of planetesimals has also been invoked (Payne et al. 2009) to explain planets on AU-scale orbits in small separation binaries.

On the other hand, there are also factors complicating planetesimal growth. In particular, eccentricity of the gaseous disk induced by the companion may affect planetary growth at small sizes ( $d \lesssim d_{gas}$ ). Also, gas drag-induced inspiral of planetesimals may deplete the disk of some solids. The relative importance of these factors for planet formation in binaries will be assessed in the future.

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